

Chiral Perturbation Theory for Vector Mesons

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We derive a heavy vector-meson chiral Lagrangian in which the vector mesons are treated as heavy static matter fields. The unknown couplings of the chiral Lagrangian are further related using the $1/N_c$ expansion. Chiral perturbation theory is applied to the vector-meson mass matrix. At one loop there are large corrections to the individual vector-meson masses, but the singlet-octet mixing angle remains almost unchanged. The parity-violating s -wave $\phi \rightarrow \rho\pi$ weak decay amplitude is derived in the combined chiral and large N_c limits. Rare ϕ decays provide a sensitive test of nonleptonic neutral current structure.

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An important application of chiral perturbation theory is to describe the interactions of matter fields (such as nucleons [1] or hadrons containing a heavy quark [2]) with low-momentum pseudo Goldstone bosons—the pions, kaons and, etc. In this paper we use chiral perturbation theory to describe the interactions of the ρ , K^* , ϕ , and ω vector mesons with low-momentum pseudo Goldstone bosons. The results of this work are relevant for experiments at the ϕ factory being built at Frascati [3]. We apply chiral perturbation theory to transitions of the form $V \rightarrow V'X$, where V and V' are vector mesons. The mass differences between the nine lowest-lying vector mesons are small compared with the chiral symmetry breaking scale of ~ 1 GeV, so chiral perturbation theory is applicable as a systematic expansion procedure for such decays. Chiral perturbation theory has previously been used to study processes such as $\rho \rightarrow \pi\pi$, which do not have a vector meson in the final state. Decays such as $\rho \rightarrow \pi\pi$ do not have soft pions in the final state, so the application of the chiral Lagrangian to such processes is not justified and should be considered as a phenomenological model.

The pseudo Goldstone boson fields can be written as a 3×3 special unitary matrix

$$\Sigma = \exp \frac{2i\Pi}{f}, \quad (1)$$

$$\mathcal{O}_\mu = \begin{bmatrix} \rho_\mu^0/\sqrt{2} + \phi_\mu^{(8)}/\sqrt{6} & \rho_\mu^+ & K_\mu^{*+} \\ \rho_\mu^- & -\rho_\mu^0/\sqrt{2} + \phi_\mu^{(8)}/\sqrt{6} & K_\mu^{*0} \\ K_\mu^{*-} & \bar{K}_\mu^{*0} & -2\phi_\mu^{(8)}/\sqrt{6} \end{bmatrix}, \quad (5)$$

and as a singlet

$$S_\mu = \phi_\mu^{(0)}. \quad (6)$$

Under chiral $SU(3)_L \times SU(3)_R$,

$$\mathcal{O}_\mu \rightarrow U \mathcal{O}_\mu U^\dagger, \quad S_\mu \rightarrow S_\mu, \quad (7)$$

where

$$\Pi = \begin{bmatrix} \pi^0/\sqrt{2} + \eta/\sqrt{6} & \pi^+ & K^+ \\ \pi^- & -\pi^0/\sqrt{2} + \eta/\sqrt{6} & K^0 \\ K^- & \bar{K}^0 & -2\eta/\sqrt{6} \end{bmatrix}. \quad (2)$$

Under chiral $SU(3)_L \times SU(3)_R$, $\Sigma \rightarrow L\Sigma R^\dagger$, where $L \in SU(3)_L$ and $R \in SU(3)_R$. At leading order in chiral perturbation theory, f can be identified with the pion or kaon decay constant ($f_\pi \simeq 132$ MeV, $f_K \simeq 160$ MeV). It is convenient when describing the interactions of the pseudo Goldstone bosons with other fields to introduce

$$\xi = \exp \frac{i\Pi}{f} = \sqrt{\Sigma}. \quad (3)$$

Under chiral $SU(3)_L \times SU(3)_R$,

$$\xi \rightarrow L\xi U^\dagger = U\xi R^\dagger, \quad (4)$$

where, in general, U is a complicated function of L, R , and the meson fields Π . For transformations $V = L = R$ in the unbroken $SU(3)_V$ subgroup, $U = V$.

The vector-meson fields are introduced as a 3×3 octet matrix

and under charge conjugation,

$$C\mathcal{O}_\mu C^{-1} = -\mathcal{O}_\mu^T, \quad CS_\mu C^{-1} = -S_\mu, \quad (8)$$

$$C\xi C^{-1} = \xi^T.$$

We construct a chiral Lagrangian for vector mesons by treating the vector mesons as heavy static fields

[4,5] with fixed four-velocity v^μ , $v^2 = 1$. The three polarization states of vector mesons with velocity v^μ satisfy $v \cdot S = v \cdot \mathcal{O} = 0$. The chiral Lagrange density which describes the interactions of the vector mesons with the low-momentum π , K , and η mesons has the general structure

$$\mathcal{L} = \mathcal{L}_{\text{kin}} + \mathcal{L}_{\text{int}} + \mathcal{L}_{\text{mass}}. \quad (9)$$

Processes such as $\phi \rightarrow K^+ K^-$ cause a net disappearance of vector mesons and can be taken into account by including anti-Hermitian terms in \mathcal{L} . The vector-meson widths are small compared with M_K and will be neglected in our analysis.

At leading order in the derivative and quark mass expansions

$$\mathcal{L}_{\text{kin}} = -i S_\mu^\dagger (v \cdot \partial) S^\mu - i \text{Tr} \mathcal{O}_\mu^\dagger (v \cdot \mathcal{D}) \mathcal{O}^\mu \quad (10)$$

and

$$\begin{aligned} \mathcal{L}_{\text{int}} = & ig_1 S_\mu^\dagger \text{Tr} (\mathcal{O}_\nu A_\lambda) v_\sigma \epsilon^{\mu\nu\lambda\sigma} + \text{H.c.} \\ & + ig_2 \text{Tr} (\{\mathcal{O}_\mu^\dagger, \mathcal{O}_\nu\} A_\lambda) v_\sigma \epsilon^{\mu\nu\lambda\sigma}, \end{aligned} \quad (11)$$

where

$$\mathcal{D}^\nu \mathcal{O}^\mu = \partial^\nu \mathcal{O}^\mu + [V^\nu, \mathcal{O}^\mu] \quad (12)$$

and

$$\begin{aligned} V^\mu = & \frac{1}{2} (\xi \partial^\mu \xi^\dagger + \xi^\dagger \partial^\mu \xi), \\ A^\mu = & \frac{i}{2} (\xi \partial^\mu \xi^\dagger - \xi^\dagger \partial^\mu \xi). \end{aligned} \quad (13)$$

The terms in \mathcal{L}_{kin} appear with minus signs because the polarization vector is spacelike. Charge conjugation invariance requires that the product of \mathcal{O}_μ^\dagger and \mathcal{O}_ν in the second term of Eq. (11) be an anticommutator. (It is important to remember that in the matrix \mathcal{O}_μ^\dagger the field $\rho_\mu^{-\dagger}$ is not equal to ρ_μ^+ , etc. In heavy vector-meson chiral perturbation theory, ρ_μ^+ destroys a ρ^+ but it does not create the corresponding antiparticle. A separate field $\rho_\mu^{-\dagger}$ is introduced to create a ρ^- .) Finally, to linear order in the quark mass expansion,

$$\begin{aligned} \mathcal{L}_{\text{mass}} = & \mu_0 S_\mu^\dagger S^\mu + \mu_8 \text{Tr} \mathcal{O}_\mu^\dagger \mathcal{O}^\mu \\ & + \lambda_1 \left[\text{Tr} (\mathcal{O}_\mu^\dagger \mathcal{M}_\xi) S^\mu + \text{H.c.} \right] \\ & + \lambda_2 \text{Tr} (\{\mathcal{O}_\mu^\dagger, \mathcal{O}^\mu\} \mathcal{M}_\xi) \\ & + \sigma_0 \text{Tr} \mathcal{M}_\xi S_\mu^\dagger S^\mu + \sigma_8 \text{Tr} \mathcal{M}_\xi \text{Tr} \mathcal{O}_\mu^\dagger \mathcal{O}^\mu, \end{aligned} \quad (14)$$

where \mathcal{M} is the quark mass matrix $\mathcal{M} = \text{diag}(m_u, m_d, m_s)$ and

$$\mathcal{M}_\xi = \frac{1}{2} (\xi \mathcal{M} \xi + \xi^\dagger \mathcal{M} \xi^\dagger). \quad (15)$$

Note that the fields S and \mathcal{O} appearing in Eqs. (10)–(12) are understood to be velocity-dependent fields which are rescaled by a common phase factor (ei-

ther $e^{-i\mu_0 v \cdot x}$ or $e^{-i\mu_8 v \cdot x}$). (The velocity-dependent vector-meson fields are related to the vector-meson fields by $\phi_\nu^\mu = \sqrt{2m} e^{imv \cdot x} \phi^\mu$ and have dimension 3/2.) This rescaling removes either μ_0 or μ_8 from Eq. (14), so only the singlet-octet mass difference $\Delta\mu \equiv \mu_0 - \mu_8$ is relevant. Phenomenologically, the parameter $\Delta\mu < 200$ MeV is comparable to mass splittings of order m_s , so in our power counting we treat $\Delta\mu$ as a quantity of order m_q . $\Delta\mu$ is of order $1/N_c$, and so vanishes in the large N_c limit.

We begin by considering the spectrum of vector mesons produced at leading order in chiral perturbation theory. The analysis is identical to the well-known SU(3) analysis [6]. Neglecting isospin breaking due to the up and down quark mass difference, i.e., $m_u = m_d = \hat{m}$, we find that

$$m_\rho = \bar{\mu}_8 + 2\lambda_2 \hat{m}, \quad m_{K^*} = \bar{\mu}_8 + \lambda_2 (\hat{m} + m_s), \quad (16a)$$

and the $\phi^{(0)} - \phi^{(8)}$ mass matrix is

$$M^{(08)} = \begin{bmatrix} \bar{\mu}_0 & -\frac{2}{\sqrt{6}} \lambda_1 (m_s - \hat{m}) \\ -\frac{2}{\sqrt{6}} \lambda_1 (m_s - \hat{m}) & \bar{\mu}_8 + \frac{2}{3} \lambda_2 (\hat{m} + 2m_s) \end{bmatrix}, \quad (16b)$$

where

$$\bar{\mu}_0 = \mu_0 + \sigma_0 \text{Tr} \mathcal{M}, \quad \bar{\mu}_8 = \mu_8 + \sigma_8 \text{Tr} \mathcal{M}. \quad (17)$$

(Isospin breaking effects have been studied in Ref. [7].) Using Eqs. (16a) and (16b), it is possible to express the elements of $M^{(08)}$ in terms of the measured vector-meson masses (up to a sign ambiguity for $M_{12}^{(08)}$)

$$M_{11}^{(08)} = m_\omega + m_\phi - \frac{4}{3} m_{K^*} + \frac{1}{3} m_\rho, \quad (18a)$$

$$M_{22}^{(08)} = \frac{4}{3} m_{K^*} - \frac{1}{3} m_\rho, \quad (18b)$$

$$\begin{aligned} M_{12}^{(08)} = M_{21}^{(08)} = & \pm \left[\left(\frac{4}{3} m_{K^*} - \frac{1}{3} m_\rho - m_\omega \right) \right. \\ & \times \left. \left(m_\phi - \frac{4}{3} m_{K^*} + \frac{1}{3} m_\rho \right) \right]^{1/2}. \end{aligned} \quad (18c)$$

The eigenstates of $M^{(08)}$ are parametrized by a mixing angle Θ_V

$$|\phi\rangle = \sin \Theta_V |\phi^{(0)}\rangle - \cos \Theta_V |\phi^{(8)}\rangle, \quad (19a)$$

$$|\omega\rangle = \cos \Theta_V |\phi^{(0)}\rangle + \sin \Theta_V |\phi^{(8)}\rangle, \quad (19b)$$

where Eqs. (18) imply the usual SU(3)_V prediction for the tangent of the mixing angle

$$\tan \Theta_V = \mp \sqrt{\frac{m_\phi - \frac{4}{3} m_{K^*} + \frac{1}{3} m_\rho}{\frac{4}{3} m_{K^*} - \frac{1}{3} m_\rho - m_\omega}} \simeq \mp 0.76. \quad (20)$$

In the large N_c limit [8,9], quark loops are suppressed so that the leading diagrams in the meson sector contain a single quark loop. As a result, the octet and singlet mesons can be combined into a single “nonet” matrix

$$N_\mu = \mathcal{O}_\mu + \frac{I}{\sqrt{3}} S_\mu, \quad (21)$$

which enters the chiral Lagrangian. The kinetic, interaction, and mass terms at leading order in $1/N_c$ are

$$\mathcal{L}_{\text{kin}} \rightarrow -i \text{Tr} N_\mu^\dagger (v \cdot \mathcal{D}) N_\mu, \quad (22)$$

$$\mathcal{L}_{\text{int}} \rightarrow i g_2 \text{Tr} (\{N_\mu^\dagger, N_\nu\} A_\lambda) v_\sigma \epsilon^{\mu\nu\lambda\sigma}, \quad (23)$$

and

$$\mathcal{L}_{\text{mass}} \rightarrow \mu \text{Tr} N_\mu^\dagger N^\mu + \lambda_2 \text{Tr} (\{N_\mu^\dagger, N^\mu\} \mathcal{M}_\xi). \quad (24)$$

Comparing with Eqs. (10)–(14), one finds that, in the $N_c \rightarrow \infty$ limit,

$$\Delta\mu \rightarrow 0, \quad \sigma_0 \rightarrow \frac{2\lambda_2}{3}, \quad \sigma_8 \rightarrow 0, \quad (25)$$

$$g_1 \rightarrow \frac{2g_2}{\sqrt{3}}, \quad \lambda_1 \rightarrow \frac{2\lambda_2}{\sqrt{3}}, \quad \tan \Theta_V \rightarrow \frac{1}{\sqrt{2}}, \quad (26)$$

the $|\phi\rangle$ state becomes “pure” $|s\bar{s}\rangle$, and the nonet matrix is

$$N_\mu = \begin{bmatrix} \rho_\mu^0/\sqrt{2} + \omega_\mu/\sqrt{2} & \rho_\mu^+ & K_\mu^{*+} \\ \rho_\mu^- & -\rho_\mu^0/\sqrt{2} + \omega_\mu/\sqrt{2} & K_\mu^{*0} \\ K_\mu^{*-} & \bar{K}_\mu^{*0} & \phi_\mu \end{bmatrix}. \quad (27)$$

If the minus sign is chosen in Eq. (18c), the prediction for the mixing angle at leading order in chiral perturbation theory, Eq. (20), is close to its value for large N_c . At leading order in chiral perturbation theory, the partial width for the Zweig forbidden decay $\phi \rightarrow \rho\pi$ summed over all three modes is

$$\Gamma(\phi \rightarrow \rho\pi) = \frac{2h^2 |\mathbf{p}_\pi|^3}{\pi f^2}. \quad (28)$$

The coupling, h , which vanishes as $N_c \rightarrow \infty$, is

$$h = \frac{g_1}{\sqrt{2}} \sin \Theta_V - \frac{g_2}{\sqrt{3}} \cos \Theta_V. \quad (29)$$

The measured branching ratio gives $h \approx 0.05$, which also suggests that the couplings are close to the $N_c \rightarrow \infty$ values.

In the nonrelativistic constituent quark model, assuming the $|\phi\rangle$ is pure $|s\bar{s}\rangle$, $g_1 = 2/\sqrt{3}$ and $g_2 = 1$. In the nonrelativistic chiral quark model [10], g_1 and g_2 are reduced by a factor of 0.75 from their values in the nonrelativistic constituent quark model.

In chiral perturbation theory the leading corrections to the expressions for the vector-meson masses in Eqs. (16) are of order $m_q^{3/2}$ (recall we are treating $\Delta\mu$ as of order m_q) and arise from one-loop self-energy diagrams giving

$$\begin{aligned} \delta m_\rho &= -\frac{1}{12\pi f^2} \left[g_2^2 \left(\frac{2}{3} m_\pi^3 + 2m_K^3 + \frac{2}{3} m_\eta^3 \right) + g_1^2 m_\pi^3 \right], \\ \delta m_{K^*} &= -\frac{1}{12\pi f^2} \left[g_2^2 \left(\frac{3}{2} m_\pi^3 + \frac{5}{3} m_K^3 + \frac{1}{6} m_\eta^3 \right) + g_1^2 m_K^3 \right], \\ \delta M_{11}^{(08)} &= -\frac{1}{12\pi f^2} g_1^2 (3m_\pi^3 + 4m_K^3 + m_\eta^3), \\ \delta M_{22}^{(08)} &= -\frac{1}{12\pi f^2} \left[g_2^2 \left(2m_\pi^3 + \frac{2}{3} m_K^3 + \frac{2}{3} m_\eta^3 \right) + g_1^2 m_\eta^3 \right], \\ \delta M_{12}^{(08)} = \delta M_{21}^{(08)} &= \frac{1}{12\pi f^2} \sqrt{\frac{2}{3}} g_1 g_2 (-3m_\pi^3 + 2m_K^3 + m_\eta^3). \end{aligned} \quad (30)$$

The singlet-octet mixing angle Θ_V including these corrections is

$$\tan \Theta_V = \mp \sqrt{\frac{m_\phi - \frac{4}{3} m_{K^*} + \frac{1}{3} m_\rho - \delta m}{\frac{4}{3} m_{K^*} - \frac{1}{3} m_\rho - m_\omega + \delta m}}, \quad (31)$$

where

$$\begin{aligned} \delta m &= -\frac{4}{3} \delta m_{K^*} + \frac{1}{3} \delta m_\rho + \delta M_{22}^{(08)} \\ &= -\frac{1}{12\pi f^2} \left(g_1^2 + \frac{2}{3} g_2^2 \right) \left(\frac{1}{3} m_\pi^3 - \frac{4}{3} m_K^3 + m_\eta^3 \right). \end{aligned} \quad (32)$$

Using the relation between g_2 and g_1 in Eq. (26), we find that

$$\delta m \rightarrow -\frac{2g_2^2}{12\pi f^2} \left(\frac{1}{3} m_\pi^3 - \frac{4}{3} m_K^3 + m_\eta^3 \right). \quad (33)$$

With $g_2 = 0.75$, Eq. (33) yields $\delta m \approx -4$ MeV. The combination of mass shifts, δm , that affects the mixing angle Θ_V is very small even though the corrections to the individual masses are substantial (e.g., $\delta m_\rho \approx -300$ MeV). δm , which is of order $1/N_c$, must transform like a **27** of flavor SU(3). The linear combination of the cubed pseudoscalar meson mass in Eq. (33) transforms like a **27**, and is numerically small. This same linear combination occurs in the violation of the Gell-Mann–Okubo formula for baryon masses [11].

For N_c large, the $\phi \rightarrow \rho\pi$ decay amplitude is of order $N_c^{-3/2}$, since the leading order $1/\sqrt{N_c}$ amplitude is forbidden by Zweig’s rule. At leading order in chiral perturbation theory it occurs at tree level because of order $1/N_c$ deviations from the relations $\tan \Theta_V = 1/\sqrt{2}$ and $g_1/g_2 = 2/\sqrt{3}$. At order $m_s \ln m_s$ in the chiral expansion the order $N_c^{-3/2}$ contribution arises from one-loop vertex and wave function corrections calculated with vertices from the nonet Lagrange density. The π and η loops do not contribute when one uses the nonet Lagrangian. The resulting decay amplitude for each of the three $\rho\pi$ modes is

$$\begin{aligned} \mathcal{A}(\phi \rightarrow \rho\pi) &= \frac{i}{f} \epsilon^{\mu\nu\lambda\sigma} \epsilon_\mu(\phi) \epsilon_\nu^*(\rho) p_{\pi\lambda} v_\sigma (2\sqrt{m_\phi m_\rho}) \\ &\times \left[\sqrt{2} h - g_2^2 \left(\frac{m_K^2}{8\pi^2 f^2} \right) \ln \left(\frac{m_K^2}{\mu^2} \right) + \dots \right], \end{aligned} \quad (34)$$

where the ellipsis denotes terms of higher order in the chiral and $1/N_c$ expansions. The terms of order m_s have a dependence on the subtraction point μ which cancels that of the logarithm in Eq. (34). With $\mu = 1$ GeV, $g_2 = 0.75$, the magnitude of the term of order $m_s \ln m_s$ in Eq. (34) is about 1.5 times as large as the measured $\phi \rightarrow \rho\pi$ decay amplitude. This suggests either that g_2 is smaller than the chiral quark model value or that there

is a partial cancellation between order $m_s \ln m_s$ and order m_s contributions to the decay amplitude.

The Frascati ϕ factory is expected to produce of order 10^{10} ϕ 's, allowing even very rare ϕ decay processes to be experimentally accessible. The $\phi \rightarrow \rho \pi$ decay amplitude has a small parity-violating s -wave amplitude that is induced by the weak interactions. This amplitude can be predicted in the combined limits of chiral $SU(2)_L \times SU(2)_R$ symmetry and large N_c . In these limits the part of the weak Hamiltonian that dominates the s -wave $\phi \rightarrow \rho \pi$ amplitude is due to Z^0 exchange,

$$\mathcal{H}_W = \eta \frac{G_F}{2\sqrt{2}} \left(1 - \frac{4}{3} \sin^2 \theta_W\right) \times (\bar{s}_\alpha \gamma_\mu s_\alpha) [\bar{u}_\beta \gamma^\mu \gamma_5 u_\beta - \bar{d}_\beta \gamma^\mu \gamma_5 d_\beta], \quad (35)$$

where $\eta \sim 1.56$ [12] arises from QCD scaling between the weak scale and low energies. In the large N_c limit, $\log \eta$ is of order $1/N_c$ times logarithms of the form $\log M_W/\Lambda$, and we have chosen to include corrections of this form [13]. The $\phi \rightarrow \rho \pi$ matrix element takes the form

$$\langle \rho \pi | \mathcal{H}_W | \phi \rangle = \eta \frac{G_F}{2\sqrt{2}} \left(1 - \frac{4}{3} \sin^2 \theta_W\right) f_\phi \epsilon_\mu(\phi) \times \langle \rho \pi | \bar{u} \gamma^\mu \gamma_5 u - \bar{d} \gamma^\mu \gamma_5 d | 0 \rangle, \quad (36)$$

where the ϕ decay constant f_ϕ is defined by

$$\langle 0 | \bar{s} \gamma_\mu s | \phi \rangle = f_\phi \epsilon_\mu(\phi). \quad (37)$$

The measured $\phi \rightarrow e^+ e^-$ decay width implies that $f_\phi \approx (492 \text{ MeV})^2$. The left-handed isovector current $\bar{u} \gamma^\mu P_L u - \bar{d} \gamma^\mu P_L d$ transforms as $(3_L, 1_R)$ under chiral $SU(2)_L \times SU(2)_R$. For matrix elements between the vacuum and a ρ plus soft pions, this current is represented by the operator

$$\bar{u} \gamma^\mu P_L u - \bar{d} \gamma^\mu P_L d = \frac{f_\rho}{2\sqrt{2}m_\rho} \text{Tr}(\xi \mathcal{O}_\mu^\dagger \xi^\dagger \tau^3), \quad (38)$$

where f_ρ is defined analogous to f_ϕ and has the value $f_\rho \approx (407 \text{ MeV})^2$ from the $\rho \rightarrow e^+ e^-$ partial width. In Eq. (38), ξ and \mathcal{O}_μ are the 2×2 matrix analogs of the corresponding 3×3 matrices used in the case of chiral $SU(3)_L \times SU(3)_R$. The right-handed isovector current is given by exchanging ξ and ξ^\dagger in Eq. (38). Using Eq. (38) to evaluate the matrix element in Eq. (36) we

find that in the combined chiral and large N_c limits

$$\begin{aligned} \langle \rho^+ \pi^- | \mathcal{H}_W | \phi \rangle_{s\text{-wave}} &= - \langle \rho^- \pi^+ | \mathcal{H}_W | \phi \rangle_{s\text{-wave}} \\ &= -i \eta \frac{G_F}{\sqrt{2}} \left(1 - \frac{4}{3} \sin^2 \theta_W\right) \\ &\quad \times \left(\frac{f_\phi f_\rho}{f}\right) \epsilon^*(\rho) \cdot \epsilon(\phi) \end{aligned} \quad (39)$$

and

$$\langle \rho^0 \pi^0 | \mathcal{H}_W | \phi \rangle_{s\text{-wave}} = 0. \quad (40)$$

Interference between the s -wave and p -wave amplitudes is possible for aligned ϕ 's, but it requires a final state interaction phase. The s -wave $\rho^+ \pi^-$ branching ratio is 10^{-11} , which is too small to be measured at the Frascati ϕ factory. However, an enhancement of the parity-violating decay rate could make the signal observable. This provides a very interesting test of new physics, because it probes nonleptonic neutral currents involving strange quarks. We will consider the application of chiral perturbation theory to other processes such as $\phi \rightarrow \rho \gamma \gamma$ elsewhere.

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- [1] J. Gasser, M.E. Sainio, and A. Švarc, Nucl. Phys. **B307**, 779 (1988); E. Jenkins and A.V. Manohar, Phys. Lett. B **255**, 558 (1991); **259**, 353 (1991).
 - [2] M.B. Wise, Phys. Rev. D **45**, 2188 (1992); G. Burdman and J.F. Donoghue, Phys. Lett. B **280**, 287 (1992); T.-M. Yan *et al.*, Phys. Rev. D **46**, 1148 (1992); P. Cho, Nucl. Phys. **B396**, 183 (1993).
 - [3] *The DAΦNE Physics Handbook*, edited by L. Maiani, G. Pancheri, and N. Paver (INFN, Frascati, Italy, 1992).
 - [4] E. Eichten and B. Hill, Phys. Lett. B **234**, 511 (1990).
 - [5] H. Georgi, Phys. Lett. B **240**, 447 (1990).
 - [6] J.J.J. Kokkedee, *The Quark Model* (Benjamin, New York, 1969).
 - [7] R. Urech, Report No. hep-ph/9504238.
 - [8] G. 't Hooft, Nucl. Phys. **B72**, 461 (1974).
 - [9] G. Veneziano, Nucl. Phys. **B117**, 519 (1976).
 - [10] A. Manohar and H. Georgi, Nucl. Phys. **B234**, 189 (1984).
 - [11] E. Jenkins, Nucl. Phys. **B368**, 190 (1992).
 - [12] J. Dai, M.J. Savage, J. Liu, and R. Springer, Phys. Lett. B **271**, 403 (1991).
 - [13] W. Bardeen, A. Buras, and J.M. Gerard, Nucl. Phys. **B293**, 787 (1987).